Correspondence of consistent and inconsistent spin-3/2 couplings via the equivalence theorem

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Abstract

Field theoretic description of higher-spin particles imposes strict consistency conditions on their interactions. On the spin 3/2 example we show that "consistent" and "inconsistent" interactions can be related by field transformations which leave observables invariant. At least in massive theories, to each "inconsistent" interaction we are able to put in correspondence a "consistent" one with the same S-matrix.

Consistent quantum field theoretic formalism for *free* higher-spin (HS) particles is fairly well established [1]. Although a covariant HS field has more spin degrees of freedom (DOF) than is necessary to describe polarizations of the physical particle, the free action can be written such that the number of DOF is reduced to the physical value. In particular, the kinetic term of HS actions has gauge symmetries which reduce the number of independent DOF down to 2, while the mass term (partially) breaks these symmetries such that the DOF number is $2s + 1$. Obviously, interactions must support this mechanism of DOF reduction, and otherwise they are *inconsistent* with the free theory construction and bound to give rise to various pathological effects, such as negative norm states and acausal modes [2,3]. Many general forms of interaction are ruled out on these grounds (see *e.g.* [4,5]).

It is straitforward to argue that the interactions which have the same type of gauge-symmetries as the kinetic term cannot change the DOF content of the free theory, and hence are consistent, see *e.g.* [6]. The problem then is to find interactions which would support the gauge-symmetries of all the HS fields involved. Unfortunately, this problem becomes highly nontrivial once "minimal" interactions to fundamental HS fields, such as that of electromagnetism and gravity, need to be included. For instance, the simplest consistent model which includes the minimal electromagnetic coupling of the spin-3/2 particle is the $N = 2$ supergravity [7], where the spin-3/2 particle can be nothing else than gravitino, so indivisible are its properties from the spacetime geometry. For many higher spin fields even such super-constrained implementations are unavailable.

In this Letter we demonstrate that an *inconsistent* interaction of a *massive* spin-3/2 field can be related to a *consistent* one by a redefinition of the spin-3/2 field. The redefinition gives also rise to some higher-order (in the coupling constant) interactions, which however cannot, in general, spoil or improve the consistency of the theory. Most interestingly, according to the equivalence theorem [8] the two theories related by the field redefinition are equivalent at the level of S-matrix elements (or, observables). Thus we obtain a recipe to overcome inconsistencies of a given interaction by including some specific higher-order interactions.

This has an important consequence for the effective field theory (EFT) formalisms, where all the necessary higher-order interactions are included anyway. In EFT's any inconsistent and a corresponding consistent coupling of massive HS field are guaranteed to be fully equivalent, as all the differences arising at the S-matrix level can be accommodated by a shift in the coefficients of some higher-order interactions. In other words, in EFT's consistency requirements generic to HS fields may be relaxed.

In the hadronic sector an EFT description, namely Chiral Perturbation Theory (ChPT), of particles with spin upto 3/2 is already extensively used, see *e.g.* [10]. Our results will in particular justify the treatment of the spin-3/2 baryons in ChPT from the viewpoint of field-theoretic consistency.

We begin with the Lagrangian of the free massive spin-3/2 Rarita-Schwinger (RS) field[1]:

$$\mathcal{L}_{\mathrm{RS}} = \overline{\psi}_\mu(x)\, \Lambda^{\mu\nu}(i\partial)\, \psi_\nu(x)\,, \quad \Lambda^{\mu\nu}(i\partial) \equiv \gamma^{\mu\nu\alpha}\, i\partial_\alpha - m\, \gamma^{\mu\nu}\,. \tag{1}$$

Corresponding field equations are

$$\Lambda^{\mu\nu}(i\partial)\, \psi_\nu = 0 = \gamma_\mu \Lambda^{\mu\nu}(i\partial)\, \psi_\nu = \partial_\mu \Lambda^{\mu\nu}(i\partial)\, \psi_\nu\,. \tag{2}$$

or, equivalently, $(i\gamma\cdot\partial - m)\psi_\mu = 0 = \gamma\cdot\psi = \partial\cdot\psi$. The kinetic term is, up to a total derivative, invariant under the gauge transformation:

$$\psi_\mu(x) \to \psi_\mu(x) + \partial_\mu \epsilon(x), \tag{3}$$

where $\epsilon(x)$ is a spinor.

The covariant propagator of the RS field in the momentum space takes the well-known form:

$$S^{\mu\nu}(p) = \frac{\not{p}+m}{p^2-m^2+i\varepsilon}\left[-\eta^{\mu\nu} + \tfrac{1}{3}\gamma^\mu\gamma^\nu + \frac{1}{3m}(\gamma^\mu p^\nu - \gamma^\nu p^\mu) + \frac{2}{3m^2}p^\mu p^\nu\right], \tag{4}$$

and satisfies

$$S^{\mu\alpha}(p)\, \Lambda^{\beta\nu}(p)\, \eta_{\alpha\beta} = \Lambda^{\mu\alpha}(p)\, S^{\beta\nu}(p)\, \eta_{\alpha\beta} = \eta^{\mu\nu}\,. \tag{5}$$

In considering the interactions, we will focus on the *linear coupling* of the spin-3/2 field, *i.e.*,

$$\mathcal{L}_{\mathrm{int}} = g\, \overline{\psi}_\mu\, j^\mu + g\, \overline{j}^\mu\, \psi_\mu\,, \tag{6}$$

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[1]Our conventions: metric tensor $\eta^{\mu\nu} = \mathrm{diag}(1,-1,-1,-1)$; γ-matrices γ^μ, $\gamma^5 = i\gamma^0\gamma^1\gamma^2\gamma^3$, $\{\gamma^\mu,\gamma^\nu\} = 2\eta^{\mu\nu}$; fully antisymmetrized products of γ-matrices $\gamma^{\mu\nu} = \frac{1}{2}[\gamma^\mu,\gamma^\nu] = \gamma^\mu\gamma^\nu - \eta^{\mu\nu}$, $\gamma^{\mu\nu\alpha} = \frac{1}{2}\{\gamma^{\mu\nu},\gamma^\alpha\} = i\varepsilon^{\mu\nu\alpha\beta}\gamma_\beta\gamma_5$, $\gamma^{\mu\nu\alpha\beta} = \frac{1}{2}[\gamma^{\mu\nu\alpha},\gamma^\beta] = i\varepsilon^{\mu\nu\alpha\beta}\gamma_5$; spinor indices are omitted.

where j can depend on fields other than ψ; g is a coupling constant. Consistency requirements would impose a condition on j^μ. For example, if the coupling is to be symmetric under the gauge transformation (3), then j must be divergenceless: $\partial \cdot j = 0$. Suppose, however, that our j does not obey any such condition and the coupling is *inconsistent*.

A field redefinition:

$$\psi_\mu(x) \to \psi_\mu(x) + g\,\xi_\mu(x)\,, \tag{7}$$

gives rise to a new linear coupling $\mathcal{L}'_{\rm int}$ plus a quadratic (in the coupling constant) interaction $\mathcal{L}_C$:

$$\begin{aligned} \mathcal{L}_{\rm RS} + \mathcal{L}_{\rm int} \to \mathcal{L}_{\rm RS} &+ \mathcal{L}'_{\rm int} + \mathcal{L}_C\,, \\ \mathcal{L}'_{\rm int} &= g\,\overline{\psi}\cdot(j + \mathit{\Lambda}\cdot\xi) + \text{H.c.}, \\ \mathcal{L}_C &= g^2\,[\overline{\xi}\cdot\mathit{\Lambda}\cdot\xi + \overline{\xi}\cdot j + \overline{j}\cdot\xi]\,. \end{aligned} \tag{8}$$

The point of this procedure is that field ξ_μ can always be chosen such that the new linear coupling $\mathcal{L}'_{\rm int}$ is *consistent*, for example,

$$\xi_\mu = (m\,\gamma^{\mu\nu})^{-1} j^\nu = -\frac{1}{m}\mathcal{O}^{(-1/3)}_{\mu\nu} j^\nu \tag{9}$$

where $\mathcal{O}^{(x)}_{\mu\nu} \equiv \eta_{\mu\nu} + x\gamma_\mu\gamma_\nu$. Then

$$j'^\mu = \gamma^{\mu\nu\alpha} i\partial_\alpha \xi_\nu \tag{10}$$

and consistency condition $\partial \cdot j' = 0$ is explicitly obeyed. Note that in this case ξ and hence $\mathcal{L}_C$ are independent of ψ. Therefore, we state that (i) an inconsistent linear coupling of a massive spin-3/2 can in general be transformed, by a redefinition of the spin-3/2 field, into a consistent coupling plus an additional quadratic coupling that does not involve the spin-3/2 field.

Furthermore, both the Lagrangian and the field transformation satisfy the conditions of the equivalence theorem [9], and thus (ii) the description in terms of $\mathcal{L}_{\rm int}$ or $\mathcal{L}'_{\rm int} + \mathcal{L}_C$ are equivalent at the level of S-matrix.

Moving the quadratic coupling to the other side of the equation, we obtain a corollary of statements (i) and (ii): given any inconsistent linear coupling we can find the supplementary second-order interaction which will provide us with the description of observables identical to the one in a consistent interacting theory.

It is interesting that the model $\mathcal{L}_{\rm RS}+\mathcal{L}_{\rm int}-\mathcal{L}_C$ will still be detected as *inconsistent*, at least by the standard analysis [2–5,11–13]. At the same time it is "S-matrix-equivalent" to the consistent model $\mathcal{L}_{\rm RS}+\mathcal{L}'_{\rm int}$. This in particular raises the question whether the HS pathologies (*e.g.*, acausal propagations) may at all manifest themselves in observables. Hopefully this question can be answered by studying the above field transformation in a constrained path-integral formulation. There one should also be able to prove the equivalence theorem for this case which is quite special because the field transformation affects the symmetries and thus the DOF content of the theory.

To demonstrate the above statements we now consider a specific example of the spin-3/2 coupling to a spin-0 and a spin-1/2 field. Such couplings are frequently used in describing the coupling of the decuplet baryons to the pion and nucleon. In particular, the conventional $\pi N\Delta$ coupling reads [11]:

$$\mathcal{L}_{\pi N\Delta} = g\overline{\psi}^i_\mu \left(\eta^{\mu\nu} + z\gamma^\mu\gamma^\nu\right) T^a_{ik}\Psi^k\partial_\nu\phi^a + \text{H.c.}, \tag{11}$$

where g is a dimensionfull coupling constant, and z is an "off-shell parameter". Herein we have retained the isospin since it plays some role in what follows. The pseudo scalar fields ϕ^a, correspond to the pion $\pi^a = (\pi^+, \pi^-, \pi^0)$; spinors Ψ^k correspond to the nucleon $N^k = (p, n)$; the RS fields ψ^i_μ represent the $\Delta^i = (\Delta^{++}, \Delta^+, \Delta^0, \Delta^-)$; T^a denotes the isospin 1/2 to 3/2 transition matrices satisfying $T^{\dagger a}T^b = \frac{2}{3}\delta^{ab} - \frac{1}{3}i\varepsilon^{abc}\tau^c$; τ^c are the isospin Pauli matrices.

This is a typical example of an "inconsistent" spin-3/2 coupling. For $z \neq -1$ it explicitly violates the constraints of the free RS theory [11], while for $z = -1$ it gives rise to the Johnson-Sudarshan and Velo-Zwanziger problems [12] (see Ref. [13] for more details and references).

To find a corresponding consistent coupling we make the field transformation (7) with

$$\xi_\mu = -\frac{1}{m}\mathcal{O}^{(-1/3)}_{\mu\varrho}\mathcal{O}^{(z)\varrho\nu}T^a\Psi\partial_\nu\phi^a \tag{12}$$

and thus obtain a consistent linear coupling [13]:

$$\mathcal{L}'_{\pi N\Delta} = -\frac{ig}{2m}\overline{G}_{\mu\nu}\gamma^{\mu\nu\lambda}T^a\Psi\,\partial_\lambda\phi^a + \text{H.c.}\,, \tag{13}$$

where $G_{\mu\nu} = \partial_\mu\psi_\nu - \partial_\nu\psi_\mu$, plus the second-order term:

$$\mathcal{L}_{\pi\pi NN} = -\left(\frac{g}{m}\right)^2 \overline{\Psi}\,\mathcal{O}^{(x)}_{\varrho\mu}\left(\gamma^{\mu\nu\alpha}\,i\partial_\alpha + m\,\gamma^{\mu\nu}\right)\mathcal{O}^{(x)}_{\nu\sigma}\,T^{\dagger b}T^c\Psi\,(\partial^\varrho\phi^{\dagger b})(\partial^\sigma\phi^c)\,. \tag{14}$$

with $x = -\frac{1}{3}(1+z)$. Thus, $\mathcal{L}_{\text{RS}} + \mathcal{L}_{\pi N\Delta} \to \mathcal{L}_{\text{RS}} + \mathcal{L}'_{\pi N\Delta} + \mathcal{L}_{\pi\pi NN}$.

To check that the field transformation leaves the S-matrix invariant let us first consider some simplest matrix elements involving the two vertices:

$$\begin{aligned} &\mathit{\Gamma}^{\mu\,a}(k) \equiv \mathit{\Gamma}^\mu(k)\,T^a, \quad \mathit{\Gamma}^\mu(k) = g\,(\eta^{\mu\nu} + z\gamma^\mu\gamma^\nu)k_\nu \quad \text{(inconsistent)} \\ &\tilde{\mathit{\Gamma}}^{\mu\,a}(k,p) \equiv \tilde{\mathit{\Gamma}}^\mu(k,p)\,T^a, \quad \tilde{\mathit{\Gamma}}^\mu(k,p) = -(g/m)\,\gamma^{\mu\nu\alpha}k_\nu p_\alpha \quad \text{(consistent)}\,. \end{aligned}$$

The Δ production amplitude is apparently the same for both vertices

$$\overline{u}(p')\,\mathit{\Gamma}^{\mu\,a}(p'-p)\,u_\mu(p) = \overline{u}(p')\,\tilde{\mathit{\Gamma}}^{\mu\,a}(p'-p,p)\,u_\mu(p) = g\,\overline{u}(p')\,(p'-p)^\mu\,u_\mu(p)\,T^a, \tag{15}$$

where $u(p)$ is the nucleon spinor, $u_\mu(p)$ is the free RS vector-spinor satisfying $(\not{p} - m)u_\mu = 0 = p\cdot u = \gamma\cdot u$, with $p^2 = m^2$.

However for the Δ-exchange amplitudes in pion-nucleon scattering, Fig. 1, the two vertices yield quite different results. The inconsistent coupling involves the spin-1/2 sector of the RS propagator, and therefore the exchange amplitude,

$$\mathit{\Gamma}^\mu(k')S_{\mu\nu}(p+k)\mathit{\Gamma}^\nu(k) \sim \frac{g^2}{m-(p+k)\cdot\gamma}P^{3/2}_{\mu\nu}(p+k)\,k'^\mu k^\nu + \text{“spin-1/2 background”}\,, \tag{16}$$

contains the controversial spin-1/2 background contributions, in addition to the spin-3/2 propagation represented by the spin-3/2 projection operator:

$$P^{3/2}_{\mu\nu}(p) = g_{\mu\nu} - \frac{1}{3}\gamma_\mu\gamma_\nu - \frac{1}{3p^2}(\not p\gamma_\mu p_\nu + p_\mu\gamma_\nu\not p). \tag{17}$$

In contrast, the consistent coupling, because of the property $p \cdot \tilde{\Gamma}(k,p) = 0$, gives rise to only the spin-3/2 propagation [13,6], namely

$$\tilde{\Gamma}^\mu(k', p_s) S_{\mu\nu}(p_s) \tilde{\Gamma}^\nu(k, p_s) = \frac{g^2}{m - \not p_s} \frac{p_s^2}{m^2} P^{3/2}_{\mu\nu}(p_s)\, k'^\mu k^\nu \,. \tag{18}$$

where $p_s = p + k = p' + k'$.

Because the decomposition into the pure spin-3/2 and spin-1/2 sector is nonlocal (the projection operators are singular at $p^2 = 0$), it is not at all obvious that the difference between the amplitudes (16) and (18) can be compensated by the local contact term $\mathcal{L}_{\pi\pi NN}$. However it indeed happens, as can most easily be seen for the case $z = -1$, when

$$\tilde{\Gamma}^\mu(k,p) - \Gamma^\mu(k) = -(g/m)\Lambda^{\mu\nu}(p)\, k_\nu \,. \tag{19}$$

Using this identity and Eq. (5) we find for the s-channel exchange the difference between the amplitudes (16) and (18) is

$$\begin{aligned} M_C^{(s)} &= [\tilde{\Gamma}(k', p_s) \cdot S(p_s) \cdot \tilde{\Gamma}(k, p_s) - \Gamma(k') \cdot S_{\mu\nu}(p_s) \cdot \Gamma(k)]\, T^{\dagger a} T^b \\ &= \Big[-(g/m)\,(k' \cdot \Gamma(k) + \Gamma(k') \cdot k) \,+ (g/m)^2 k' \cdot \Lambda(p_s) \cdot k\Big]\; T^{\dagger a} T^b. \end{aligned} \tag{20}$$

A similar contribution comes from the u-channel graph such that the total difference sums up into

$$M_C^{(s)} + M_C^{(u)} = -(g/m)^2\, (T^{\dagger a} T^b - T^{\dagger b} T^a)[\tfrac{1}{2}\gamma^{\mu\nu\alpha}\, (p + p')_\alpha + m\, \gamma^{\mu\nu}] k'_\mu k_\nu \,, \tag{21}$$

which is exactly canceled by the contact interaction $\mathcal{L}_{\pi\pi NN}$.

Clearly, Green's functions which do not represent observable quantities need not be the same. For instance, the one-loop Δ self-energy will be different for the two couplings. However, at the level of the S-matrix the equivalence is restored. That is, the amplitude containing the self-energy with the consistent coupling (l.h.s. in Fig. 2) is identical to the one-loop amplitude with the inconsistent coupling plus the contact term (r.h.s. in Fig. 2). It is easy to convince oneself that this equivalence will persist to any number of loops.

The higher-order contact term may be absent in some cases. For instance, in the case of $\pi N\Delta$ coupling with *neutral pions* only, obtained by neglecting the isospin complications and taking the real scalar field in Eq. (11), the contact term $\mathcal{L}_{\pi\pi NN}$ vanishes. In this case the inconsistent (conventional) and the consistent (RS gauge-invariant) $\pi N\Delta$ couplings are one-to-one equivalent at the S-matrix level.

Recently there have been similar findings specifically concerning the conventional $\pi N\Delta$ coupling in πN scattering. Tang and Ellis showed [14] that the contribution of the "off-shell parameter" in the Δ-exchange amplitude can be absorbed into a contact term, and therefore, they argued, this parameter is redundant in ChPT. In Ref. [15] it was shown numerically that, once contact terms of the form of ρ and σ meson exchanges are included, the gauge-invariant Eq. (13) and conventional Eq. (11) couplings at the tree level give the same prediction for the πN threshold parameters provided some coupling constants are readjusted. Here we have basically proven the results of these observations in a general fashion, extending to any S-matrix elements and the quantum level.

With respect to the linear couplings in general, we can argue that within an EFT framework any linear spin-3/2 coupling is acceptable, in the sense that, even if it is inconsistent (by standard criteria), it gives a description equivalent to a consistent coupling. That is because in EFT the additional $\mathcal{L}_C$ type of terms, which provide the equivalence of the inconsistent and consistent couplings, are to be included anyway with arbitrary coefficients and in both situations. Thus, if the effective Lagrangian with an inconsistent coupling has $\mathcal{L}_C$ term with arbitrary coefficient c_1, it is S-matrix-equivalent to the Lagrangian with a consistent coupling and $\mathcal{L}_C$ term with a different but yet arbitrary coefficient $c_2 = c_2(c_1, g/m)$. In other words, all the differences are completely accounted for by a change in the coupling constant(s).

Nevertheless, let us also emphasize that the use of consistent (gauge-invariant) couplings makes the calculations much easier and more transparent. In particular, the spin-1/2 sector can be entirely dropped from the RS propagator [6,13], while analyzing the spin-3/2 self-energy, one does not need to consider the ten scalar functions of the most general tensor structure [16], but only two of them [15,17], just as in the spin-1/2 case. Besides the technical advantages, consistent couplings involve the physical higher-spin contributions only, and hence they are preferable in the analysis of separate contributions and effects due to spin-3/2 particles versus the rest. This can be important when the properties of separate HS resonances, such as the $\Delta(1232)$-isobar, are being extracted in a model-dependent way from experimental data, see *e.g.* Refs. [18,19,15].

The case of couplings *quadratic* in the spin-3/2 field includes the fundamental problem of the inconsistencies in the theory of charged spin-3/2 particle [2,3,7,20]. But much of the said about the linear couplings is applicable to this case as well. Given any inconsistent coupling of a massive RS field ψ_μ, we can obtain an on-shell equivalent consistent coupling by the replacement:

$$\psi_\mu \to \frac{i}{m} \mathcal{O}^{(-1/3)}_{\mu\lambda}\, \gamma^{\lambda\alpha\beta}\, \partial_\alpha \psi_\beta \,. \tag{22}$$

It is then possible to work out the exact field transformation relating the couplings and the supplementary higher-order terms providing their equivalence at the S-matrix level. Working these out for any specific example is beyond the scope of this letter. Without going into details it is already clear that such transformation must be nonlinear in the RS field and the number of supplementary terms is infinite.

In conclusion, we have shown how by making a field redefinition in an *inconsistent* model of an interacting spin-3/2 field one can obtain an S-matrix-equivalent *consistent* model, and vice versa. At least for massive fields it is always possible to do so, and we can exploit that to justify the use of any inconsistent linear couplings in EFT's. The fact that an inconsistent model (with inconsistent coupling plus, if necessary, the higher-order terms) leads to the same S-matrix as a consistent model, implies that the inconsistencies generic to HS fields, *e.g.* negative-norm states and acausal propagations, do not manifest themselves in observables in that case. On the other hand, we emphasize that consistent interactions have anyway conceptual and technical advantages. Apart from being quantizable by standard methods, they do not couple to lower-spin components of HS fields hence making the calculations simpler with more transparent interpretation.

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FIGURES

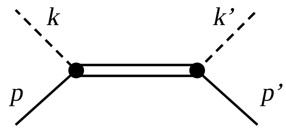

FIG. 1. The s-channel Δ-exchange in πN scattering. The dashed, solid, and double lines denote the pion, the nucleon, and the Δ, respectively.

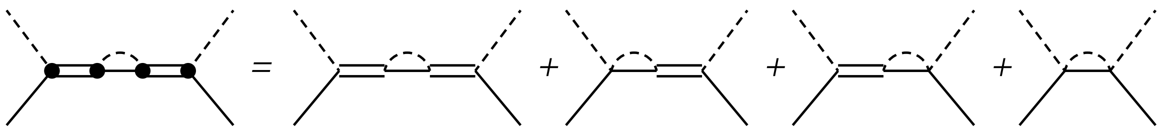

FIG. 2. The equivalence of a πN-scattering loop amplitude with the consistent $\pi N \Delta$ coupling (denoted by a dot) on the l.h.s., and in the corresponding inconsistent model on the r.h.s.; crossed graphs are omitted.